

750 GeV diphoton resonance in a visible heavy QCD axion model

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Abstract

In this paper, we revisit a visible heavy QCD axion model in light of the recent reports on the 750 GeV diphoton resonance by the ATLAS and CMS experiments. In this model, the axion is made heavy with the help of the mirror copied sector of the Standard Model while the successful Peccei-Quinn mechanism is kept intact. We identify the 750 GeV resonance as the scalar boson associated with spontaneous breaking of the Peccei-Quinn symmetry which mainly decays into a pair of the axions. We find that the mixings between the axion and η and η' play important roles in its decays and the resultant branching ratio into two photons. The axion decay length can be suitable for explaining the diphoton excess by the di-axion production when its decay constant $f_a \simeq 1$ TeV. We also find that our model allows multiple sets of the extra fermions without causing the domain wall problem, which is advantageous to explain the diphoton signal.

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I. INTRODUCTION

The success of Kabayashi-Maskawa mechanism [1] for CP violation in the quark sector as well as the Sakharov conditions for baryogenesis in the early universe [2, 3] strongly suggest that CP violation is an intrinsic structure of nature. If this is indeed the case, it naturally leads to the question why the strong interaction conserves the CP symmetry so well when it is allowed to have its own CP -violating parameter, the θ angle. This is the strong CP problem in QCD.

The Peccei-Quinn mechanism [4, 5] is the most attractive solution to the strong CP problem. As a prominent prediction, this mechanism comes with a pseudo-Nambu-Goldstone boson, the axion, of mass of $O(100)$ keV [6, 7]. Such a light axion, however, has been excluded by extensive experimental searches [8].

To circumvent the experimental search constraints, there are two approaches: one way is to make the axion couplings to the standard-model particles very weak, and the other is to make the axion heavier. The former lead to the well-known invisible axion models [9–12] in which the axion decay constant f_a is taken to be very large, *e.g.*, $f_a > 10^9$ GeV, so that the axion couplings are highly suppressed. The invisible axion models have, however, a serious drawback. A global U(1) symmetry is an essential ingredient of the Peccei-Quinn mechanism. However, such a global symmetry is most likely broken explicitly by gravitational effects. These effects emerge as higher dimensional operators suppressed by the Planck scale in the effective field theory which shift the θ angle. By remembering the stringent upper bound on the θ angle, $\theta \lesssim 10^{-10}$ [8], the shift in the θ angle by the gravitational effects is unacceptably large even for the original axion model [6, 7] where the decay constant is the electroweak scale. The situation becomes even worse for invisible axion models with $f_a \gtrsim 10^9$ GeV.

In parallel to the invisible axion models, many people have also tried to construct models involving the Peccei-Quinn mechanism with a heavy axion, which turns out to be very difficult [13]. Among various attempts, however, an exceptionally successful idea was proposed by Rubakov [14] where a mirror world of the Standard Model was introduced. Recently, a concrete and viable model of such a heavy axion model has been constructed [15] (see also Refs. [16–18]¹) and is called a visible heavy QCD axion model. In that work, we showed

¹ The models discussed in Refs. [16, 17] have various unsolved cosmological problems.

that the axion decay constant could be as low as $f_a \simeq O(1)$ TeV and axion mass was around $m_a > O(0.1)$ GeV without any conflict with experimental, astrophysical or cosmological constraints. This model is visible in the sense that it predicts a scalar boson and vector-like fermions at $O(1)$ TeV in addition to the heavy axion. It should be emphasized that the above-mentioned gravitational breaking effects are sufficiently small, thanks to such a small f_a and the heaviness of the axion.

In this paper, we revisit a visible heavy QCD axion model in light of the recent reports on the 750 GeV diphoton resonance by the ATLAS and CMS experiments [19, 20]. As a result, we show that the scalar boson predicted in the visible axion model with $f_a \simeq 1$ TeV can be identified with mass 750 GeV and mainly decays into two axions. Each axion decays into two photons with a sizable branching ratio, so that the di-axion signal mimics the diphoton signal. So far, many works have been done to discuss the excess and some of them [21–24] (see also [25–28] for related works) try to explain the signal using axion-like particles. Our model is the first realistic model which explains the diphoton excess by using the QCD axion that solves the strong CP problem.

The paper is organized as follows. In section II, we briefly summarize the status of the heavy axion as well as the model of Ref. [15]. In section III, we discuss how the model explains the diphoton resonance by identifying it as the scalar boson in the visible axion model. We carefully discuss how the axion decays since the scalar boson mainly decays into a pair of the axions. Finally, discussions and conclusions are given in section IV.

II. HEAVY QCD AXION MODEL

In order to achieve a heavy axion, Rubakov [14] proposed an idea to employ a mirror copy of the Standard Model. By assuming a \mathbb{Z}_2 symmetry between the Standard Model and its mirror copy, the θ -angles in these two sectors are aligned at the high energy input scale, such as the Planck scale. The electroweak symmetry breaking scale and the dynamical scale of QCD in the copied sector can be varied once the \mathbb{Z}_2 symmetry is spontaneously broken while the θ angles are intact. With the higher scales in the copied sector, the axion mass is enhanced without spoiling the success of the Peccei-Quinn mechanism.

In Ref. [15], Harigaya and three of the authors propose a concrete realization of this mechanism. We prepare the Standard Model with a single Higgs doublet and its mirror

copy. To implement the Peccei-Quinn mechanism, we introduce extra colored vector-like fermions ψ and ψ' in each sector and identify their chiral symmetries as the U(1) Peccei-Quinn symmetry. Hereafter, objects with a prime (') refer to those in the copied sector. A complex scalar ϕ is introduced to break the Peccei-Quinn symmetry spontaneously. As in the KSVZ axion model [9, 10], ϕ couples to ψ and ψ' via

$$\Delta\mathcal{L} = g\phi (\bar{\psi}\psi + \bar{\psi}'\psi') , \quad (1)$$

where g denotes the coupling constant. Assuming that the Peccei-Quinn symmetry is spontaneously broken by the vacuum expectation value (VEV) of ϕ , we decompose ϕ into an axion a and a scalar boson s around its VEV, $f_a/\sqrt{2}$,

$$\phi = \frac{1}{\sqrt{2}}(f_a + s)e^{ia/f_a} . \quad (2)$$

In terms of the axion, the Peccei-Quinn symmetry is realized non-linearly by

$$a/f_a \rightarrow a/f_a + \alpha , \quad \alpha \in [0, 2\pi) . \quad (3)$$

By using a \mathbb{Z}_2 breaking spurion σ , the squared mass parameters of Higgs and Higgs' can be varied with each other [15]. Besides, the dynamical scale of QCD (Λ) and QCD' (Λ') can be also varied by introducing some new scalar particles in both sectors whose masses are different from each other by the spurion effect. By making the electroweak' scale and Λ' larger than those in the Standard Model sector, the axion mass is dominated by the contributions from the copied sector:

$$m_a \propto \frac{f_{\pi'} m'_{\pi}}{f_a} , \quad (4)$$

which can be much larger than conventional models. Here, m'_{π} and f'_{π} are the mass and the decay constant of the pion' in the copied sector.

In Ref. [15], we examined astrophysical and cosmological constraints on the heavy axion model. We found that the axion heavier than the QCD phase transition temperature, $T_{\text{QCD}} = \mathcal{O}(100) \text{ MeV}$, satisfied the astrophysical/cosmological constraints for at least $f_a \lesssim 10^5 \text{ GeV}$. In fact, with such a heavy axion, all the copied sector particles decouple from the thermal bath of the Standard Model above the QCD phase transition, and hence, the contributions of the copied sector to the effective number of relativistic species are diluted

(see also discussions at the end of section III B).²

Several remarks on the cosmological constraints are in order. When we assume the seesaw mechanism [29] in both sectors, the neutrino masses in the copied sector are proportional to the Higgs' VEV squared, leading to too much (hot/warm) dark matter as the VEV of Higgs' is about hundreds times larger than the VEV of the Higgs. To avoid this problem, it is safer to assume that the seesaw mechanism does not take place in the copied sector³, so that the neutrino' decays into pion' and charged lepton' in the copied sector. The nucleons', electron' and pions' in the copied sector are stable, which places upper limits on the electroweak' and the QCD' scales in the copied sector to avoid too much dark matter abundance [15]. It is also important to arrange the model so that ψ and ψ' mix with the quarks and quarks' so that they decay fast enough not to cause cosmological problems.

The constraints from the rare meson decays as well as the beam dump experiments do not exclude a heavy axion for $m_a \gtrsim 3 \times m_\pi$. In this mass range, the axion decay length is short since it has $a \rightarrow 3\pi$ decay modes (see next section), and hence, it decays well before reaching the detectors of the beam dump experiments such as the CHARM experiment [30]. Besides, due to the lack of direct axion-quark couplings as in the KSVZ type models, the rate for a rare meson decay into an axion is suppressed by the mixing between the axion and the neutral pion. As a result, the constraints from the rare meson decays are also insignificant in this mass region.

Let us also consider the constraint on the model from the LHC. The searches for extra quarks at the LHC put lower limits on the mass of an additional fermion, $gf_a/\sqrt{2}$ (see Eq. (1)). As alluded to before, ψ decays into quarks via the ψ -quark mixing. The experimental lower bounds on the masses of extra quarks are 800–900 GeV [31–34], depending on the branching ratios of ψ into b and t quarks. Assuming $g \sim 1$ if we require it not to blow up below the Planck scale, we find $f_a \gtrsim 1$ TeV.

Finally, let us discuss how the Peccei-Quinn mechanism is durable to explicit breaking of the Peccei-Quinn symmetry by the Planck-suppressed operators. At the leading order, the

² The copied sector is in thermal equilibrium with the Standard Model particles even after the QCD' phase transition.

³ This is possible by switching off the $(B - L)'$ symmetry breaking in the copied sector by using the \mathbb{Z}_2 spurion.

explicit breaking terms are given by dimension-five operators,

$$\mathcal{L} = \frac{\kappa}{5!M_{\text{PL}}} (\phi^5 + \phi^{*5}) , \quad (5)$$

with a coefficient κ . Such higher-dimensional operators lead to a non-vanishing θ angle at the minimum of the axion potential:

$$\Delta\theta_{\text{eff}} \sim 10^{-11} \times \kappa \left(\frac{f_a}{10^3 \text{ GeV}} \right)^3 \left(\frac{1 \text{ GeV}}{m_a} \right)^2 , \quad (6)$$

which is consistent with the current upper limit for $f_a \simeq 1 \text{ TeV}$ and $m_a = O(1) \text{ GeV}$. This feature is quite favorable compared with invisible axion models where the Peccei-Quinn mechanism can be easily spoiled by explicit breaking terms suppressed by the Planck scale.

III. INTERPRETATION OF THE EXCESS

A. Properties of the scalar resonance

In this section, we discuss whether the reported 750 GeV diphoton resonance can be identified with the scalar boson s , the radial component of ϕ (see Eq. (2)). For that purpose, let us first discuss the decay width of s .

Assuming that ψ is sufficiently heavy, the effective interactions involving s are given by

$$\begin{aligned} \Delta\mathcal{L} = & \frac{s}{f_a} \partial_\mu a \partial^\mu a + \frac{\alpha_s}{8\pi} \frac{g}{\sqrt{2}M_D} f(t_D) s G^{\mu\nu} G_{\mu\nu} + \frac{\alpha_2}{8\pi} \frac{g}{\sqrt{2}M_L} f(t_L) s W^{\mu\nu} W_{\mu\nu} \\ & + \frac{2\alpha_Y}{8\pi} \left(\frac{1}{3} \frac{g}{\sqrt{2}M_D} f(t_D) + \frac{1}{2} \frac{g}{\sqrt{2}M_L} f(t_L) \right) s B^{\mu\nu} B_{\mu\nu} + (G, B, W \rightarrow G', B', W') , \end{aligned} \quad (7)$$

where we assume that the copied couplings, $\alpha'_s, \alpha'_2, \alpha'_Y$, are the same as ones of the Standard Model at UV due to the \mathbb{Z}_2 symmetry and only slightly different at IR. For a while, we assume that the extra fermions $\psi^{(\prime)}$ and $\bar{\psi}^{(\prime)}$ form respectively the $\mathbf{5}$ and $\mathbf{5}^*$ representations of $SU(5)_{\text{GUT}}$, and name the colored fermion and doublet fermion, $\psi^{(\prime)}_{D,L}$, respectively. The parameters with the subscripts D, L are for $\psi^{(\prime)}_{D,L}$, respectively. The three gauge field strengths in each sector are denoted by $G^{(\prime)}$, $W^{(\prime)}$, and $B^{(\prime)}$, respectively. The function $f(t)$ is defined by⁴

$$f(t) \equiv t \left[1 + (1-t) \arcsin^2 \left(\frac{1}{\sqrt{t}} \right) \right] , \quad (8)$$

⁴ In the limit of heavy fermion masses, $f(t)$ converges: $\lim_{t \rightarrow \infty} f(t) \rightarrow 2/3$.

with t 's being

$$t_{D,L} = \frac{4M_{D,L}^2}{M_s^2} . \quad (9)$$

The mass of s , M_s , is taken to be 750 GeV.

As is clear from this effective Lagrangian, we find that s mainly decays into a pair of the axions, while the modes into gauge bosons are loop-suppressed. As a result, the total decay width of s is roughly given by

$$\Gamma(s \rightarrow 2a) \simeq \frac{M_s^3}{32\pi f_a^2} \simeq 4.2 \text{ GeV} \times \left(\frac{1 \text{ TeV}}{f_a} \right)^2 \quad (10)$$

for $M_s = 750 \text{ GeV}$. It should be noted that it is difficult to explain the broad width $\Gamma \simeq 45 \text{ GeV}$ that is slightly favored by the ATLAS experiment unless the decay constant $f_a \lesssim 300 \text{ GeV}$. With such a small decay constant, however, the extra colored particles are predicted to be too light.

At a first glance, it seems that this dilaton s cannot account for the diphoton signal since the 2γ decay mode of s is highly suppressed, $\text{BR}(s \rightarrow \gamma\gamma) \sim (\alpha/4\pi)^2 \sim 10^{-8}$. However, the $2a$ decay mode can mimic the diphoton signal since the axions in the final state are highly boosted and decay into collimated photons. Such photon jet events have been studied in Refs. [35–41], and for recent works [23–28] discuss this possibility in the context of 750 GeV diphoton excess.

At the LHC, the resonance s is produced via the gluon fusion process. By using the narrow width approximation, the production cross section of s is estimated to be

$$\sigma(g + g \rightarrow s) \simeq 11 \text{ fb} \times \left(\frac{f(t_D)}{2/3} \right)^2 \left(\frac{1 \text{ TeV}}{f_a} \right)^2 , \quad (11)$$

where we use the MSTW2008 parton distribution functions [42].⁵

Therefore, to account for the observed excess

$$4 \text{ fb} \lesssim \sigma \times \text{BR}(s \rightarrow 2a) \times \text{BR}(a \rightarrow 2\gamma)^2 \lesssim 10 \text{ fb} , \quad (12)$$

the axion branching ratio into photons should be of $\mathcal{O}(1)$. As will be shown below, $\text{BR}(a \rightarrow 2\gamma) = \mathcal{O}(1)$ is achieved through the mixings between the axion and the η and η' mesons in the Standard Model.⁶

⁵ See, *e.g.*, Ref. [43] for a discussion on higher-order QCD corrections.

⁶ To avoid confusion, we reserve the name η' to denotes the pseudoscalar meson in the Standard Model, and we use $(\eta)'$ and $(\eta')'$ to denote the pseudoscalar mesons in the copied sector.

B. Properties of the axion

To explain the diphoton excess by mimicking a single photon with collimated photons from the boosted axion decay, the axion needs to have $BR(a \rightarrow 2\gamma) = \mathcal{O}(1)$. For a very heavy axion higher than the QCD scale, $m_a \gg \Lambda$, the effective couplings of the axion to the Standard Model particles are given by

$$\mathcal{L}_a = \frac{\alpha_s}{8\pi} \frac{a}{f_a} G\tilde{G} + \frac{\alpha_2}{8\pi} \frac{a}{f_a} W\tilde{W} + 2 \left(\frac{1}{3} + \frac{1}{2} \right) \frac{\alpha_Y}{8\pi} \frac{a}{f_a} B\tilde{B}, \quad (13)$$

where we take the limit of $M_{D,L} \gg m_a$. In this case, the main decay mode of the axion is QCD jets, and the branching ratio into 2γ is highly suppressed:

$$BR(a \rightarrow 2\gamma) \propto (a_Y/a_s)^2 = \mathcal{O}(10^{-2}). \quad (14)$$

Thus, in order to have a sizable branching ratio into the photon mode, the axion mass is required to be within $\mathcal{O}(1)$ GeV so that the hadronic decay modes are suppressed. In the following, we examine the branching ratios of the axion for $m_a \lesssim 1$ GeV, where the axion mixings to the η and η' in the Standard Model play important roles.

To discuss how the axion mixes with η and η' in the Standard Model sector, let us consider the effective Lagrangian

$$\begin{aligned} \mathcal{L} = & \frac{1}{2} \partial_\mu a \partial^\mu a + \frac{1}{2} \partial_\mu \eta_8 \partial^\mu \eta_8 + \frac{1}{2} \partial_\mu \eta_0 \partial^\mu \eta_0 \\ & - \frac{1}{2} m_a^2 a^2 - \frac{1}{2} m_8^2 \eta_8^2 - \Delta^2 \eta_8 \eta_0 - \frac{1}{2} m_0^2 \left(\eta_0 + \frac{f_0}{\sqrt{6} f_a} a \right)^2, \end{aligned} \quad (15)$$

where η_8 and η_0 denote the neutral component of the octet and the singlet pseudo-Nambu-Goldstone modes resulting from chiral symmetry breaking of QCD, respectively. We neglect the contributions from π^0 because the mixing to π^0 is more suppressed. The parameter f_0 denotes the decay constant of η_0 , and will be fixed to $f_0 \simeq f_\pi \simeq 93$ MeV shortly. At the leading order of the chiral perturbation theory, the mass parameters m_8^2 and Δ^2 are generated from the quark mass terms (see, *e.g.*, Ref. [44])

$$m_8^2 \simeq \frac{1}{3} \frac{(m_u + m_d + 4m_s)}{(m_u + m_d)} m_\pi^2, \quad (16)$$

$$\Delta^2 \simeq \frac{\sqrt{2}}{3} \frac{(m_u + m_d - 2m_s)}{(m_u + m_d)} m_\pi^2, \quad (17)$$

while m_0^2 is generated by the anomalous breaking of $U(1)_A$.⁷

⁷ Here, we neglect the mass term of η_0 generated from the quark mass terms.

The mass term proportional to m_0^2 reflects the fact that the combination of the chiral rotation and the Peccei-Quinn symmetry,

$$\frac{\eta_0}{f_0} \rightarrow \frac{\eta_0}{f_0} + \alpha, \quad \frac{a}{f_a} \rightarrow \frac{a}{f_a} - \sqrt{6}\alpha, \quad (\alpha \in [0, 2\pi)) , \quad (18)$$

is free from the anomaly of QCD.⁸ It should be also noted that in the KSVZ models, the axion appears only through the m_0^2 term, and there is no kinetic mixing between the axion and the pseudo-Nambu-Goldstone modes.

Since the chiral symmetries and the Peccei-Quinn symmetry have the anomalies of QED, $\eta_{0,8}$ and the axion have anomalous couplings to the photons:

$$\mathcal{L} = \frac{\eta_8}{f_8} \frac{\alpha_{\text{QED}}}{4\sqrt{3}\pi} F\tilde{F} + \frac{\eta_0}{f_0} \frac{\alpha_{\text{QED}}}{\sqrt{6}\pi} F\tilde{F} + \frac{8}{3} \frac{a}{f_a} \frac{\alpha_{\text{QED}}}{8\pi} F\tilde{F}, \quad (20)$$

where f_8 denotes the decay constant of η_8 . The anomalous coupling of the axion comes from the anomalous couplings in Eq. (13).

Now, let us first resolve the mass mixing between η_8 and η_0 to obtain the mass eigenstates, η and η' . Unfortunately, it is known that the observed decay widths of η and η' into two photons cannot be reproduced by using the mass mixing parameter Δ^2 in Eq. (17).⁹ Thus, we follow instead the phenomenological approaches taken in Refs. [45, 46], in which the mixing angle between η - η' and f_0 are taken to reproduce the observed values using the anomalous couplings in Eq. (20). According to Refs. [45, 46], we find

$$\sin \theta \simeq -1/3, \quad f_0 \simeq 1 \times f_\pi, \quad (21)$$

where the mixing angle θ is defined by

$$\eta_8 = \cos \theta \eta + \sin \theta \eta', \quad (22)$$

$$\eta_0 = -\sin \theta \eta + \cos \theta \eta'. \quad (23)$$

Next, let us resolve the axion mixing with η and η' . From the mass term in Eq. (15), we

⁸ We take the normalization of η_0 such that the $U(1)_A$ symmetry corresponds to

$$\frac{\eta_0}{f_0} \rightarrow \frac{\eta_0}{f_0} + \alpha \quad \longleftrightarrow \quad q_{L,R} \rightarrow (e^{i\alpha/\sqrt{6}}, e^{i\alpha/\sqrt{6}}, e^{i\alpha/\sqrt{6}}) \times q_{L,R}, \quad (19)$$

where $q_{L,R} = (u_{L,R}, d_{L,R}, s_{L,R})$.

⁹ From the calculation of the chiral perturbation theory, f_8 is fixed to be about $1.3 \times f_\pi$ [45, 46].

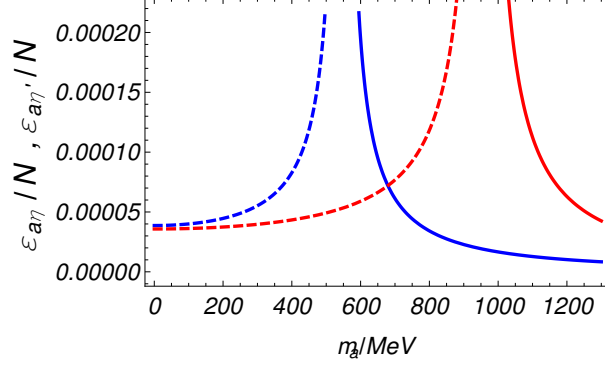


FIG. 1: The mixing angles $\varepsilon_{a\eta}$ (blue) and $\varepsilon_{a\eta'}$ (red) as a function of m_a for $f_a = 1$ TeV. The dashed curves indicate the negative mixing angles. Here, we take $\sin \theta \simeq -1/3$ and $f_0 \simeq f_\pi$. Here and hereafter, N denotes the number of colored Fermions, which we will discuss later.

find the mass eigenstates (η_D, η'_D, a_D) ;

$$\eta \simeq \eta_D + \varepsilon_{a\eta} a_D, \quad (24)$$

$$\eta' \simeq \eta'_D + \varepsilon_{a\eta'} a_D, \quad (25)$$

$$a \simeq a_D - \varepsilon_{a\eta} \eta_D - \varepsilon_{a\eta'} \eta'_D, \quad (26)$$

where

$$\varepsilon_{a\eta} = -\frac{f_0}{\sqrt{6}f_a} \frac{m_{\eta'}^2}{m_\eta^2 - m_a^2} \sin \theta, \quad \varepsilon_{a\eta'} = \frac{f_0}{\sqrt{6}f_a} \frac{m_{\eta'}^2}{m_{\eta'}^2 - m_a^2} \cos \theta. \quad (27)$$

Due to the above mixing, a coupling constant, $c_{a\mathcal{O}}$, of a_D to an operator \mathcal{O} is given by,

$$c_{a\mathcal{O}} = \varepsilon_{a\eta'} c_{\eta'\mathcal{O}} + \varepsilon_{a\eta} c_{\eta\mathcal{O}}. \quad (28)$$

For example, the axion decay width receives important contributions from the anomalous couplings of η and η' .

In Fig. 1, we show the mixing angles of the axion to η and η' as a function of m_a for $f_a = 1$ TeV. Here, we take $\sin \theta \simeq -1/3$ and $f_0 \simeq f_\pi$. The figure shows that the mixing angles are enhanced for either $m_a \simeq m_\eta$ or $m_a \simeq m_{\eta'}$. The signs of the mixing angles are opposite in the region of $m_\eta < m_a < m_{\eta'}$.

By taking into account the mixing effects, we calculate the decay widths of the relevant modes, as shown in Fig. 2. Two figures show that the widths are enhanced for either $m_a \simeq m_\eta$ or $m_a \simeq m_{\eta'}$, where the mixing to η or η' is enhanced. The figure also shows that the 2γ mode

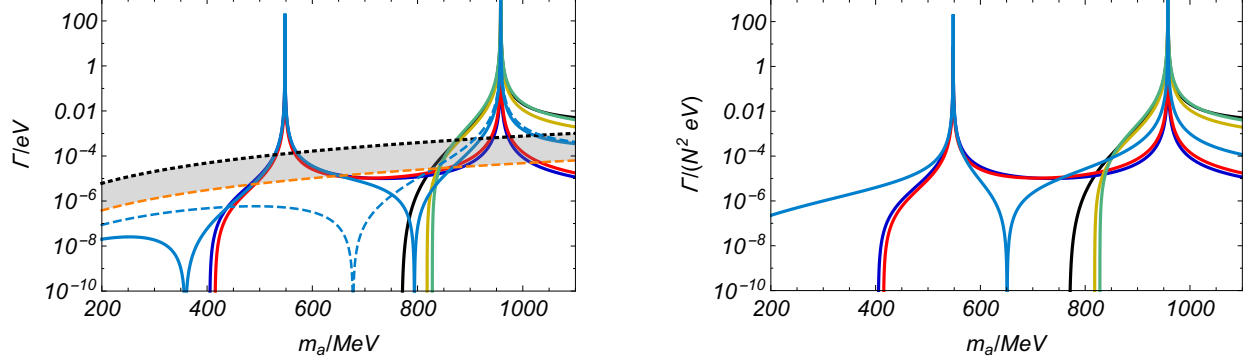


FIG. 2: (Left) The relevant partial decay widths of the axion into $\gamma\gamma$ (light blue), $3\pi^0$ (dark blue), $\pi^0\pi^+\pi^-$ (red), $\rho + \gamma$ (black), $\eta + 2\pi^0$ (grass green), $\eta + \pi^+ + \pi^-$ (green), and $\gamma'\gamma'$ (black dashed and gray band) for $f_a = 1 \text{ TeV}$. Here we take the η - η' mixing angle to have $\sin\theta \simeq -1/3$. The blue dashed line indicates $\Gamma(a \rightarrow \gamma\gamma)$ without the a - η mixing. Here we can see that the dip around 650 MeV comes from the phase cancellation. The orange dashed curve represents the decay width of the axion into $\gamma\gamma$ only from the intrinsic axion anomalous coupling to γ in Eq. (20). In this figure, we take $N = 1$. (Right) The partial decay widths for vanishing intrinsic axion anomalous couplings to γ and γ' .

is suppressed at around $m_a \simeq 670 \text{ MeV}$, which is due to destructive interference between the contributions of η and η' . The analysis of each mode is given in the appendix.

Several comments are in order. As discussed in the appendix, we estimate the 3π and $\eta+2\pi$ modes from the observed decay widths of η/η' into 3π and η' into $\eta+2\pi$.¹⁰ We approximate the decay amplitudes squared by the decay widths divided by the phase space volume. We then combine them according to Eq. (28), assuming that they add up constructively. Accordingly, our estimates of the widths into 3π and $\eta + 2\pi$ have $\mathcal{O}(1)$ uncertainties.¹¹

In the figure, we also show the decay width of the $2\gamma'$ mode as a gray band. Similar to the width of the 2γ mode, this mode also receives contributions from the mixing of the axion to $(\eta)'$, $(\eta')'$ and CP -odd glueball' in the copied sector, depending on the mass parameters. In this paper, we parametrize the anomalous coupling to the γ' by

$$\mathcal{L} \simeq \left(\frac{8}{3} + A_{\text{eff}} \right) \frac{a}{f_a} \frac{\alpha_{\text{QED}'}}{8\pi} F' \tilde{F}', \quad (29)$$

¹⁰ It should be noted that it is difficult to reproduce the observed decay widths using the chiral perturbation theory (see, for example, Ref. [47]).

¹¹ In fact, the observed Dalitz plots of $\eta \rightarrow \pi^0\pi^+\pi^-$ and $\eta' \rightarrow 3\pi^0$ are not flat [48].

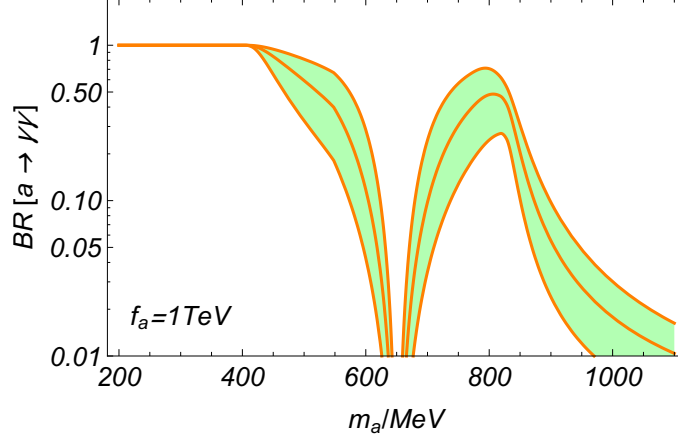


FIG. 3: The branching ratio of the axion into $\gamma\gamma$ for $f_a = 1 \text{ TeV}$. We set the η - η' mixing angle to be $\sin\theta \simeq -1/3$. The central curve in the band corresponds to $BR(a \rightarrow \gamma\gamma)$ when we estimate the widths into 3π and $\eta + 2\pi$ by using the approximation given in the appendix, while the upper and the lower curves correspond to the width multiplied by a factor of $1/3$ and 3 , respectively. We assume that the $a \rightarrow \gamma'\gamma'$ width is zero due to the anomaly cancellation.

where $\alpha_{\text{QED}'} \simeq \alpha_{\text{QED}}$ by assumption. It should be noted that A_{eff} is highly suppressed when $m_{u',d'} \gg \Lambda'$ and the mass of the axion mainly comes from the mixing to the CP -odd glueball' in the copied sector. The maximal value of $A_{\text{eff}} \simeq 8$ which is the case when the axion mainly mixes with $(\eta'_0)'$ and there is no mixing between $(\pi^0)'$ and $(\eta_8)'$.

As we discussed in [15], γ' is required to decouple from the thermal bath of the Standard Model before the hadron decoupling, which requires that the axion mass is much higher than the QCD phase transition temperature. Since the mass range of our interest is not very much larger than the QCD phase transition scale, the decoupling of γ' could have not completed before the QCD phase transition, since the Boltzmann suppression of the axion production is not significant in this mass range. In fact, for $m_a \lesssim 1 \text{ GeV}$, we find that the decoupling of γ' completes at a temperature slightly below the QCD phase transition temperature when the axion anomalous coupling to the γ' is given by the one in Eq.(29). To evade the dark radiation constraint, $N_{\text{eff}} = 3.15 \pm 0.23$ [49], we hereafter assume $A_{\text{eff}} \ll 1$ and also introduce extra matter fields in both sectors which couple to ϕ so that the intrinsic anomalous coupling is canceled. Hence, N , the total number of colored Fermions is need not to be unity. Then, the axion-gluon couplings in Eqs. (15) and (20) as well as the gluon

couplings of s in Eq. (7) are multiplied by N . We will mention the domain wall problem associated to non-unity N later. Accordingly, the intrinsic axion anomalous coupling to γ (*i.e.*, the last term in Eq. (20)) is also vanishing. In the right panel of Fig. 2, we show the branching ratios assuming vanishing intrinsic axion anomalous couplings to both γ and γ' . We also take $A_{\text{eff}} \ll 1$ assuming $m_{u',d'} \gg \Lambda'$ in the copied sector. The figure shows that vanishing anomalous coupling case is rather preferred in the sense of enough branching ratio.

C. 750 GeV di-axion resonance

As discussed above, the axion has sizable decay widths into 2γ and $3\pi^0$. When the axion is highly boosted, these modes seem to mimic the diphoton signal since each π^0 immediately decays into 2γ . However, the boosted $3\pi^0$ is not acceptable as a single photon. This is because when the number of photons is large, some of them are converted before the electromagnetic calorimeter. And once their energies are measured, it is difficult for them to fake a single photon signal¹². It should be also noted that it is not obvious even whether a collimated 2γ event is acceptable as a single photon. Considering the detector inner structure [52, 53], we conclude that the 2γ mode is safe and does not reduce the acceptance of the signal because their energies are high and the converted e^+e^- are not bent too much for the tracker to measure the energy. A more detailed study will be given elsewhere.

With these cautions in mind, we simply assume that the acceptance of $a \rightarrow 2\gamma$ as a single photon is of $\mathcal{O}(1)$, while $a \rightarrow 3\pi^0 \rightarrow 6\gamma$ is not acceptable. In Fig. 3, we show the branching ratio of the axion into 2γ for $f_a = 1$ TeV. Here, to take account of the $\mathcal{O}(1)$ uncertainties in the estimation of the widths into $a \rightarrow 3\pi$ and $a \rightarrow \eta + 2\pi$, we adopt the decay widths of those modes scaled from the ones given in Fig. 3 by a factor of 1/3 (top), 1 (middle), and 3 (bottom) for each line. The figure shows that the branching ratio of $a \rightarrow 2\gamma$ can be of $\mathcal{O}(1)$ for $400 \text{ MeV} \lesssim m_a \lesssim 600 \text{ MeV}$ and $m_a \simeq 800 \text{ MeV}$.

It is noted that the decay length of the axion should be sufficiently short to account for

¹² For a single photon, the cumulative conversion rate before it reaches to the electromagnetic calorimeter is about 20% even in the central region of the ATLAS detector (*i.e.*, $\eta = 0$) [50]. Accordingly, the conversion probability of a photon-jet consisting of six photons is naively expected to be 100% for a promptly decaying axion, where the charged tracks made by the conversion carry only a fractional energy of the axion. For a detailed analysis on the discrimination of photon-jets from single photons at the LHC detectors, see [51].

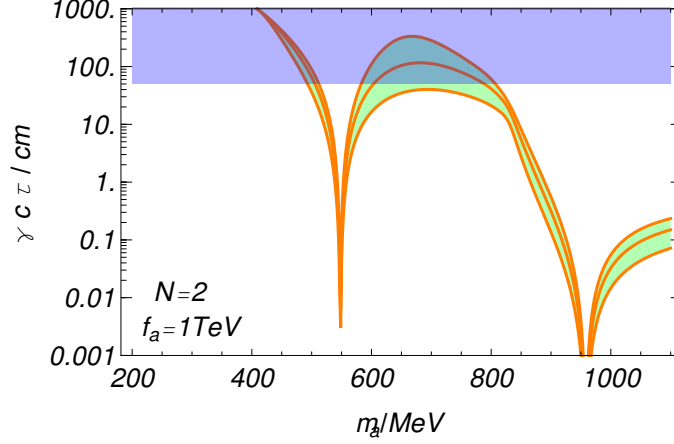


FIG. 4: The boosted decay length of an axion for $f_a = 1 \text{ TeV}$. We assume that the axion is produced from a two-body decay of a resonance with mass 750 GeV and $N = 2$. Hence, the boost factor $\gamma \simeq 325 \text{ GeV}/m_a$. The shaded blue region indicates that the boosted decay width exceeds 50 cm .

the signal. In Fig. 4, we show the decay length of an axion generated from the decay of a 750 GeV resonance for $f_a = 1 \text{ TeV}$ and $N = 2$. The figure shows that the axion decays well before the electromagnetic calorimeter. Thus, the di-axion signal in which each of the axions decays into a pair of photons can indeed mimic the diphoton signal. When the boosted decay length is longer than about 50 cm , it means that some of the axions reach the electromagnetic calorimeter before they decay.

By putting all the above discussions together, we now find the preferred range of the decay constant f_a to explain the diphoton excess by the di-axion signal with each of the axions decaying into two photons. In Fig. 5, we show the decay constant that satisfies Eq. (12) for $N = 2$. To take account of the $\mathcal{O}(1)$ uncertainties in the estimation of the widths into $a \rightarrow 3\pi$ and $a \rightarrow \eta + 2\pi$, we adopt the decay widths of these modes scaled by a factor of $1/3$ (orange), 1 (blue) and 3 (green) from the one given in Fig. 3 for each band. The figure shows that the signal can be explained for $f_a \simeq 1 \text{ TeV}$ and m_a around 500 MeV with $N \gtrsim 2$. It should be noted that the boosted decay length of the axion is longer than 50 cm when the axion is too light.

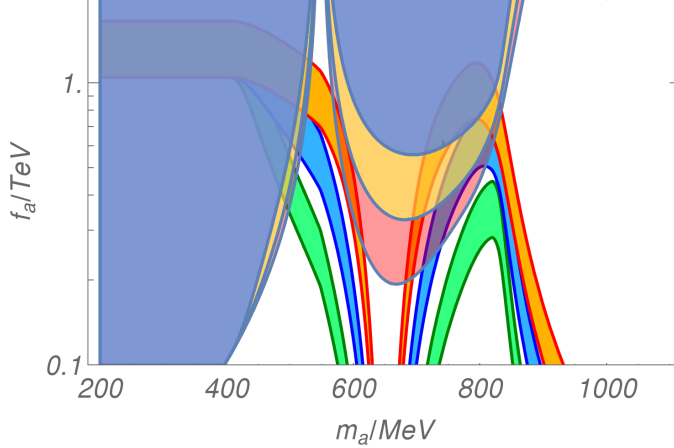


FIG. 5: The favored decay constant to reproduce the diphoton excess via the di-axion signal for $N = 2$. Each colored band corresponds to the decay constant that gives the appropriate cross section times branching ratios in Eq. (12). For each band, we adopt the decay widths of $a \rightarrow 3\pi$'s which is scaled by a factor of $1/3$ (orange), 1 (blue) and 3 (green), respectively. The shaded regions are excluded from the condition, $c\tau\gamma > 50$ cm. The red, yellow and blue regions corresponds where $\Gamma(a \rightarrow 3\pi)$ is multiplied by $1/3$, 1 and 3 , respectively.

D. Model with multiple extra fermions

As we have discussed, we may introduce multiple matters, *i.e.* non-zero N . One caveat here is the serious domain wall problem, as the discrete \mathbb{Z}_N subgroup of the Peccei-Quinn symmetry remains unbroken by the anomalies of QCD^(\prime).¹³ The discrete \mathbb{Z}_N symmetry is eventually broken spontaneously by the VEV of the axion once it attains a potential from QCD ^{\prime} effects, and the domain wall is formed when the spontaneous symmetry breakdown takes place. If the \mathbb{Z}_N symmetry is an exact symmetry, the domain walls are stable and immediately dominate the energy density of the universe.

As we have discussed at the end of section II, the visible heavy axion model is successful even if the Peccei-Quinn symmetry is explicitly broken by Planck-suppressed operators. In the presence of such extra breaking, the discrete \mathbb{Z}_N symmetry is no longer exact, since the

¹³ If the signs of Peccei-Quinn charges of the multiple extra fermions are different, it is possible to set the domain wall number to be unity even for $N > 1$.

explicit breaking generates the differences in the energy densities of vacua at the order of

$$\Delta V \sim \frac{f_a^5}{M_{\text{PL}}} , \quad (30)$$

and hence the formed domain walls are unstable [54, 55]. In fact, the domain walls get accelerated towards the domain with a higher vacuum energy density

$$A \simeq \frac{\Delta V}{T_{\text{wall}} \Lambda'^4} \simeq \frac{\Delta V}{f_a \Lambda'^2} \simeq \frac{f_a^4}{M_{\text{PL}} \Lambda'^2} . \quad (31)$$

Here, we approximate the thickness of a domain wall to be $T_{\text{wall}} \simeq f_a / \Lambda'^2$. With such an acceleration much larger than the Hubble parameter at the time of domain wall formation, the domain walls collapse immediately after their formation.¹⁴ Therefore, the models with multiple sets of fermions do not cause the domain wall problem in this model, thanks to a small axion decay constant.

Finally, let us emphasize that the extra colored fermions in multiples of three are important to evade the dark radiation constraint. As we have commented in the previous section, it is safe to consider a model where the intrinsic axion anomalous couplings to γ and γ' as in Eq. (20) and Eq. (29) are vanishing. For that purpose, we need to introduce extra colorless fermions with the opposite Peccei-Quinn charge to cancel the anomalous couplings to γ and γ' . The gauge charges of the extra fermions are restricted to allow small mixings to the Standard Model^(f) fermions to avoid stable colored/charged particles. As a result, the intrinsic anomalous couplings to γ and γ' can be cancelled by the colorless contributions only for the extra colored fermions in multiples of three.

IV. CONCLUSIONS

In this paper, we have revisited a visible heavy QCD axion model, studied in Ref. [15], in light of the recent reports on the 750 GeV diphoton resonance by the ATLAS and CMS experiments [19, 20]. In this model, the axion is made heavy with the help of the mirror copied sector of the Standard Model, thereby evading all the astrophysical/cosmological/experimental constraints even for $f_a = \mathcal{O}(1)$ TeV while preserving the successful Peccei-Quinn mechanism. The smallness of the decay constant is highly advantageous when we consider explicit breaking effects of the Peccei-Quinn symmetry by quantum gravity.

¹⁴ The domain walls mainly collapse into axions. The energy density of the emitted axions are much smaller than that in the thermal bath.

We identify the 750 GeV resonance as the scalar boson associated with spontaneous breakdown of the Peccei-Quinn symmetry whose primary decay mode is a pair of the axions. By carefully examining the decay properties of the axion, we find that its branching ratio into two photons and the boosted decay length can be suitable for explaining the observed diphoton excess using the di-axion signal for $f_a \simeq 1$ TeV. We also argue that our model allows multiple sets of extra fermions without inducing the domain wall problem. Such additional fermions are not only advantageous to explaining the signal but also important to constructing a model that evades cosmological constraints.

As seen in Fig. 4, the axion mainly decays outside the beam pipe of the LHC experiments. Since the axion has sizable branching ratios into 3π 's, these modes are expected to leave striking signals inside the detectors, with which we can test the model.

Finally, let us comment on some possibilities to achieve a broad width $\Gamma \simeq 45$ GeV that is slightly favored by the ATLAS experiment. As we have discussed in Eq. (10), such a broad width requires a small decay constant $f_a \lesssim 300$ GeV, which predicts too light extra colored particles for $g \lesssim 1$. To avoid this problem, we need to assume a strong coupling between ϕ and extra fermions, *i.e.* $g \gg 1$. Such a strong coupling blows up immediately above the TeV scale, and hence, we need some UV completion where ϕ and/or ψ 's are composite particles. Such a possibility will be discussed elsewhere.

Another way to achieve a broad width is to introduce extra fermions which are not changed under the gauge symmetries of the Standard Model nor the ones of the copied sector, so that ϕ mainly decays into those invisible extra fermions.¹⁵ In this case, we also need to assume a rather large N to make the production cross section of s enhanced to explain the diphoton excess since the branching ratio into $2a$ is small when ϕ mainly decays into the invisible extra fermions (see *e.g.* [56, 57] for related discussions).

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¹⁵ Those extra neutral fermion can be a dark matter candidate.

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Appendix A: Axion Decay Width

In this appendix, we summarize the decay width formulas for the axion.

1. $a \rightarrow \gamma\gamma$

According to Eq. (28), the axion coupling to two photon is given by

$$\mathcal{L} = C_{a\gamma\gamma} \frac{a}{f_a} \frac{\alpha_{\text{QED}}}{8\pi} F\tilde{F} , \quad (\text{A1})$$

with an effective coefficient

$$C_{a\gamma\gamma} = \frac{8}{3} + \frac{f_a}{f_\pi} \varepsilon_{a\eta'} C_{\eta'\gamma\gamma} + \frac{f_a}{f_\pi} \varepsilon_{a\eta} C_{\eta\gamma\gamma} . \quad (\text{A2})$$

Here, we define

$$C_{\eta\gamma\gamma} = \frac{2}{\sqrt{3}} \left(\frac{f_\pi}{f_{\eta^0}} \cos \theta - \sqrt{8} \frac{f_\pi}{f_{\pi_9}} \sin \theta \right) , \quad (\text{A3})$$

$$C_{\eta'\gamma\gamma} = \frac{2}{\sqrt{3}} \left(\frac{f_\pi}{f_{\eta^0}} \sin \theta + \sqrt{8} \frac{f_\pi}{f_{\pi_9}} \cos \theta \right) . \quad (\text{A4})$$

With the effective coefficient, the axion decay width is given by

$$\Gamma = \frac{1}{4\pi} \left(\frac{\alpha_{\text{QED}}}{8\pi} \right)^2 C_{a\gamma\gamma}^2 \frac{m_a^3}{f_a^2} . \quad (\text{A5})$$

2. $a \rightarrow \rho + \gamma$

To estimate the $a \rightarrow \rho + \gamma$ decay, we parameterize the $\rho + \gamma$ couplings to η and η' by effective interactions (see, *e.g.*, Ref. [58]),

$$\mathcal{L} = \frac{1}{2} \frac{\eta_8}{f} c_{\eta_8 \rho} F_\rho \tilde{F} + \frac{1}{2} \frac{\eta_0}{f} c_{\eta_0 \rho} F_\rho \tilde{F} + (F_\rho \leftrightarrow F) . \quad (\text{A6})$$

From the observed decay widths [8]

$$\Gamma_{\rho \rightarrow \eta + \gamma} \simeq 44 \text{ keV} , \quad \Gamma_{\eta' \rightarrow \rho + \gamma} = 57 \text{ keV} , \quad (\text{A7})$$

we obtain

$$\begin{aligned} c_{\eta_8 \rho} &= 9.2 \times 10^{-3} , \\ c_{\eta_0 \rho} &= 1.3 \times 10^{-2} . \end{aligned} \quad (\text{A8})$$

By suitably combining the coefficients, we derive

$$\mathcal{L} = \frac{1}{2} \frac{a^0}{f_a} c_{a\rho} F_\rho \tilde{F} + (F_\rho \leftrightarrow F) , \quad (\text{A9})$$

with

$$c_{a\rho} \simeq \varepsilon_{a\eta} \left(\frac{f_a}{f_8} c_{\eta_8 \rho} \cos \theta - \frac{f_a}{f_0} c_{\eta_0 \rho} \sin \theta \right) + \varepsilon_{a\eta'} \left(\frac{f_a}{f_8} c_{\eta_8 \rho} \sin \theta + \frac{f_a}{f_0} c_{\eta_0 \rho} \cos \theta \right) . \quad (\text{A10})$$

With this coefficient, we have

$$\Gamma = \frac{1}{2\pi} \frac{c_{a\rho}^2 (m_a^2 - M_\rho^2)^3}{f_a^2 m_a^3} . \quad (\text{A11})$$

3. $a \rightarrow 3\pi, a \rightarrow \eta + 2\pi$

Since the decay widths of η and η' are not well reproduced by the chiral perturbation theory, we estimate the axion decay width by using an approximate amplitude of η and η' into 3π ,

$$|\mathcal{A}(a \rightarrow 3\pi)| = |\varepsilon_{a\eta} \bar{\mathcal{A}}(\eta \rightarrow 3\pi^0)| + |\varepsilon_{a\eta'} \bar{\mathcal{A}}(\eta' \rightarrow 3\pi^0)| , \quad (\text{A12})$$

$$|\mathcal{A}(a \rightarrow \eta + 2\pi)| = |\varepsilon_{a\eta'} \bar{\mathcal{A}}(\eta' \rightarrow \eta + 2\pi^0)| . \quad (\text{A13})$$

Here, $|\bar{\mathcal{A}}|$'s on the right hand side are estimated by dividing the decay widths [8] by the corresponding phase space volumes and taking the square root ¹⁶:

$$|\bar{\mathcal{A}}(\eta \rightarrow \pi^0 \pi^+ \pi^-)| \simeq 0.26, \quad |\bar{\mathcal{A}}(\eta \rightarrow 3\pi^0)| \simeq 0.30, \quad (\text{A14})$$

$$|\bar{\mathcal{A}}(\eta' \rightarrow \pi^0 \pi^+ \pi^-)| \simeq 0.15, \quad |\bar{\mathcal{A}}(\eta' \rightarrow 3\pi^0)| \simeq 0.11, \quad (\text{A15})$$

$$|\bar{\mathcal{A}}(\eta' \rightarrow \eta \pi^+ \pi^-)| \simeq 6.7, \quad |\bar{\mathcal{A}}(\eta' \rightarrow \eta \pi^0 \pi^0)| \simeq 4.5. \quad (\text{A16})$$

Since we do not know how the η and η' modes interfere with each other, we simply assume constructive interference. Accordingly, our estimations of the decay widths for $a \rightarrow 3\pi$ and $a \rightarrow \eta + 2\pi$ have $\mathcal{O}(1)$ uncertainties.

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¹⁶ The amplitudes come from scalar 4-point interactions and are dimensionless.

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